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NRL Memorandum Report 6228

Development of Sidebands in Tapered and in Untapered Free-Electron Lasers

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July 1, 1988

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AD-A197 828

SECURITY CLASSIFICATION OF THIS PAGE

REPORT DOCUMENTATION PAGE

1a REPORT SECURITY CLASSIFICATION UNCLASSIFIED		1b RESTRICTIVE MARKINGS	
2a SECURITY CLASSIFICATION AUTHORITY		3 DISTRIBUTION AVAILABILITY OF REPORT Approved for public release; distribution unlimited.	
2b DECLASSIFICATION/DOWNGRADING SCHEDULE		5 MONITORING ORGANIZATION REPORT NUMBER(S)	
4 PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 6228		7a NAME OF MONITORING ORGANIZATION	
6a NAME OF PERFORMING ORGANIZATION Naval Research Laboratory	6b OFFICE SYMBOL (if applicable) Code 4790	7b ADDRESS (City, State, and ZIP Code)	
8a ADDRESS (City, State, and ZIP Code) Washington, DC 20375-5000		9 PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER	
8b NAME OF MONITORING ORGANIZATION NRL	8c OFFICE SYMBOL (if applicable)	10 SOURCE OF FUNDING NUMBERS	
8d ADDRESS (City, State, and ZIP Code) Washington, DC		PROGRAM ELEMENT NO 63221C	PROJECT NO W31RPD- 7-D4039
11 SECURITY CLASSIFICATION Development of Sidebands in Tapered and in Untapered Free-Electron Lasers			
12 PERSONAL AUTHOR(S) Bartizel, B. L., Ling, R. A., Sprangle, P. and Tang, C.M.			
13a TYPE OF REPORT Technical	13b TIME COVERED FROM TO	14 DATE OF REPORT (Year, Month, Day) 1988 July 1	15 PAGE COUNT 51
17 PERFORMING ORGANIZATION Naval Research Laboratory, Intl. Corp., McLean, VA 22102 Naval Research Laboratory, Inc., Springfield, VA 22150			
18a SUBJECT TERMS FREE-ELECTRON LASERS SUBGROUP		18b SUBJECT TERMS (Continue on reverse if necessary and identify by block number) Free-electron laser Tapered wiggler Sideband frequencies Nonadiabatic electron orbits	
<p>Abstract (Continue on reverse if necessary and identify by block number)</p> <p>A time-dependent, axisymmetric code is employed to examine the development of sidebands in a free-electron laser. For the case where the input signal undergoes an extended period of exponential growth, a broad spectrum of sidebands with growth rates comparable to that of the signal is excited. In general, in an untapered system the optical field displays considerable modulation after several synchrotron periods. An analytical model, in qualitative agreement with a number of features of the simulations, is discussed. In a tapered system the amplitude of the sidebands approaches a quasi-steady level that is several orders of magnitude below that of the untapered case, and the output optical field displays only a slight modulation. The optimal rate of tapering employed, to maximize efficiency, leads to substantial reduction in the growth rate of sidebands. This result is discussed in connection with the nonadiabatic nature of particle motion in the tapered system.</p>			
21 ABSTRACT SECURITY CLASSIFICATION UNCLASSIFIED		22a TELEPHONE (Include Area Code) (202) 767-3493	
22b OFFICE SYMBOL Code 4790			

DD FORM 173, 1-85

81 and edition may be used until exhausted
All other editions are obsolete

SECURITY CLASSIFICATION OF THIS PAGE

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DEVELOPMENT OF SIDEBANDS IN TAPERED AND IN UNTAPERED FREE-ELECTRON LASERS

I. Introduction

In a free-electron laser (FEL) the synchrotron oscillation of electrons trapped in the ponderomotive potential well may couple energy into sideband frequencies. The ensuing instability leads to the modulation of the output signal and, in consequence, to an increase in its spectral width.

The growth of sideband frequencies has been the subject of discussion in a number of papers.¹⁻⁷ In the work presented herein this process is examined by means of a time-dependent code in an assumed axisymmetric geometry. Both tapered and untapered wigglers are examined. Two regimes of sideband development are examined in detail. In one, the carrier amplitude is small and all frequencies within the linear gain spectrum develop independently. In the other, the initial amplitude of the carrier is large and hence coupled to the sideband modes via the synchrotron oscillation of the electrons. Simplified analytical models of sideband growth in the two regimes are presented as an aid to understanding the important features of the simulations.

It should be emphasized at the outset that the numerical work reported herein is intended to contrast the development of sidebands in untapered and in tapered FEL systems. In this connection, the important issue of sideband start-up from the noise spectrum appropriate to the electron beam and wiggler parameters employed is not addressed.

Manuscript approved February 16, 1988.

II. Numerical Model

This section presents the equations that form the basis for the investigation of the FEL interaction.

The electrons move under the influence of the ponderomotive force due to the beating of the wiggler field with the optical field. The computations are simplified by employing the Gaussian-Laguerre Source-Dependent Expansion (SDE) technique,^{8,9} thereby minimizing the number of Laguerre polynomials required for an accurate description of the optical field.

The optical field is taken to be of the form

$$\underline{a}_r(r, z, t) = \frac{1}{2} \underline{a}(r, z, t) \exp \left[i \left(\frac{\omega}{c} z - \omega t \right) \right] \underline{e}_x + \text{c.c.} ,$$

where $\underline{A}_r = mc^2 \underline{a}_r / |e|$ is the radiation vector potential, m is the rest-mass of an electron, $|e|$ is the magnitude of the electronic charge, c is the speed of light in vacuo, ω is the radian frequency, and \underline{e}_x is the unit vector along the x axis. The wiggler field is assumed to be plane-polarized, of amplitude B_w and period $2\pi/k_w$:

$$\underline{B}_w(z) = \frac{-i}{2} B_w \exp (ik_w z) \underline{e}_y + \text{c.c.} ,$$

where transverse variations of the wiggler field are neglected, and \underline{e}_y is the unit vector along the y axis. The equations of motion of the j -th electron, of energy $\gamma_j mc^2$, are then given by

$$\frac{d\gamma_j}{dt} = \frac{i\omega_w f_B}{4\gamma_j} \sum_n a_n L_n \left(2r_j^2 / r_s^2 \right) \exp \left[i\psi_j - (1-i\alpha) r_j^2 / r_s^2 \right] + \text{c.c.} , \quad (1)$$

$$\frac{d\psi_j}{dt} = ck_w \left(1 - \gamma_r^2 / \gamma_j^2 \right) , \quad (2)$$

where

$$\gamma_r^2 = \frac{1}{2} \left(\omega / ck_w \right) \left(1 + a_w^2 / 2 \right) \quad (3)$$

defines the resonant (i.e., synchronous) relativistic factor, $a_w = |e|B_w / mc^2 k_w$ is the normalized vector potential of the wiggler, r_j is the radial distance of the j -th electron, $\psi_j = (\omega/c + k_w) z_j - \omega t$ is the relative phase with z_j the axial location of the j -th electron, and $f_E = J_0(\xi) - J_1(\xi)$ is the usual difference of Bessel functions, with $\xi = (a_w/2)^2 / (1 + a_w^2/2)$. The results to be presented in this paper pertain to the case where only a_w is tapered and the period of the wiggler is taken to be constant.

Following Sprangle et al.,⁸ the envelope of the radiation field is expanded as follows

$$a(r, z, t) = \sum_n a_n(z, t) L_n \left[2r^2 / r_s^2 \right] \exp \left[-(1 - i\alpha) r^2 / r_s^2 \right].$$

Here, $\alpha(z, t)$ is related to the curvature of the optical wavefronts, $r_s(z, t)$ is the spot size, and $L_n (2r^2 / r_s^2)$ is the Laguerre polynomial of degree n . The method of SDE then permits a complete specification of the optical field by solving⁸

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z} \right) r_s = \frac{2c^2 \alpha}{\omega r_s} - r_s c B_I, \quad (4)$$

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z} \right) \alpha = \frac{2c^2 (1 + \alpha^2)}{\omega r_s^2} + 2c \left(B_R - \alpha B_I \right), \quad (5)$$

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z} + c A_n \right) a_n = i n B c a_{n-1} + i (n+1) B^* c a_{n+1} - i c F_n, \quad (6)$$

where

$$F_n = \frac{-v}{N\omega c} \left(\frac{2c}{r_s}\right)^2 \sum_j \left(\frac{f_B a_w}{\gamma}\right)_j L_n \left(2r_j^2/r_s^2\right) \exp \left[-i\psi_j - (1+i\alpha)r_j^2/r_s^2\right], \quad (7)$$

$$A_n = \frac{2ic}{\omega r_s} (2n + 1 - i\alpha) + i (2nB_R + B),$$

and

$$B = F_1/a_0, \\ \equiv B_R + i B_I. \quad (8)$$

In Eq. (7) the sum on j runs over the electrons in a given ponderomotive bucket and N denotes the number of electrons initially therein, and $v = I_b/(mc^2 v_z/|e|)$ is the Budker parameter, where I_b is the electron beam current.

A detailed presentation of the SDE approach is given in Ref. 8. For orientation, however, it should be noted that in vacuo one has the well-known result: $r_s(z) = r_s(0) (1+z^2/z_R^2)^{1/2}$, $\alpha = z/z_R$ where $r_s(0)$ is the minimum spot size (at $z = 0$) and $z_R = (\omega/2c) r_s^2(0)$ is the Rayleigh range. These results follow from Eqs. (4) and (5) upon neglecting B , i.e., neglecting the electron beam.

Typically 10 optical modes ($n = 0, \dots, 9$) are included in the computations. As noted in the Introduction, the question of the proper spectrum and noise level for the sidebands due to spontaneous emission is not considered herein (see Ref. 10). The seed for the sideband frequencies may be incorporated in several ways. As an example, the noise level may be estimated from the Larmor formula and spread uniformly but with random phase over the computational spectrum. Although there are fluctuations between runs with different initial random phases, the general trend of

sideband development is as described herein. Finally, in all the computations the initial electron distribution is taken to be monoenergetic, the radial profile of the electron beam is taken to be parabolic, and betatron motion is neglected.

III. Numerical Results

For definiteness the parameters for the computations presented herein correspond to those of the Paladin experiment at the Lawrence Livermore National Laboratory,¹¹ and are listed in Table I.

In the linear regime the maximum growth rate obtains at the resonance wavelength $\lambda_{\text{res}} = 10.34 \text{ } \mu\text{m}$. In what follows, where appropriate, a wavelength λ will be denoted in terms of the relative shift $\delta \equiv (\lambda/\lambda_{\text{res}}) - 1$ from the resonant value. Note that the wavelength of the carrier is given by $2\pi c/\omega$. Spectra are obtained by performing a spatial Fourier series analysis of the optical field $a(r=0, z, t)$ along the mesh. The result, displayed in the figures as Fourier amplitude, is dimensionless.

a) Multi-frequency, small input power

Simulations of a pulse of radiation extending over many ponderomotive buckets permit the development of frequencies besides that of the input signal, and as electrons slip relative to the radiation pulse the sideband instability may develop.

Figure 1 shows the development of 3 spectral components of the optical field when the input signal is at $10.6 \text{ } \mu\text{m}$ ($\delta = 2.51 \times 10^{-2}$) and 100 W. Of course there are many other spectral components besides those shown in Fig. 1; however, the curves shown do indicate the general trend in the development of the sideband frequencies. Note in particular that in the exponential regime the growth rate of the component at $\delta = 0.77 \times 10^{-2}$ exceeds that of the main signal. This result is discussed in Section IV.

b) Multi frequency, untapered magnet

Figure 2 shows the development of the carrier for the case where 300 MW of $10.6 \text{ } \mu\text{m}$ radiation ($\delta = 2.51 \times 10^{-2}$) is injected into an untapered magnet. Again, there are many modes in the spectrum that grow along the length of the wiggler. The dashed curve in Fig. 2 indicates the maximum

amplitude, or the envelope, of the rest of spectrum as a function of z . Fig. 2 also shows in detail the evolution of one of the fastest growing modes, at $\delta = 5.71 \times 10^{-2}$, indicating the trend in the development of the instability. At the end of the wiggler (25 m) the amplitude of the sidebands is large enough to spatially modulate the optical field by about 30%. Figure 3, which shows the phase of the carrier along the wiggler, will be discussed in connection with the analytical model for the instability in Section IV. At the wiggler exit the electron beam distribution function has a clear multi-stream character as indicated in Fig. 4. The spectrum of the optical field at this point consists principally of the carrier with a group of Stokes and anti-Stokes modes on either side, Fig. 5. The approximately symmetrical form of the spectrum, which is a reflection of comparable growth rates for modes symmetrically disposed with respect to the carrier, is discussed in Section IV.

c) Multi-frequency, tapered magnet

As is well known, for practical purposes the magnet employed in an FEL device must be tapered so as to enhance its efficiency and extraction. Figure 6 shows the development of the carrier and of the maximum amplitude of all the sidebands through a device where the normalized vector potential $a_w = |e|B_w/mc^2k_w$ is tapered as shown in Fig. 7. The form of the tapering employed in the computations is obtained by simply prescribing a constant rate of decrease of energy for a synchronous electron, at an assigned radius. From Eq. (3) with $dy_r/dz = \text{constant}$, one obtains $a_w(z)$. Comparing Figs. 2 and 6 it is apparent that in the tapered device the sidebands originate at a level substantially below that of the carrier. As expected, the optical field is observed to be only slightly modulated in space. The electron distribution function (not shown here) consists principally of two

groups, an untrapped group and a trapped group at lower energy. Finally, Fig. 7 also shows the efficiency of the FEL, with 16% being a ten-fold improvement over the peak efficiency for the untapered device.

IV. Analysis of Results

An understanding of Fig. 1 may be obtained by performing a single-bucket linear stability analysis of Eqs. (1)-(7). The presentation is limited to the fundamental optical mode, and a monochromatic electron beam of energy $\gamma_0 mc^2$ per electron.

Defining

$$\dot{\psi}_0 = ck_w (1 - \gamma_r^2 / \gamma_0^2),$$

$$\tilde{\psi}_j = \psi_j - \dot{\psi}_0 t,$$

$$\Gamma_j = \gamma_j / \gamma_0,$$

$$A = ia \exp(i\dot{\psi}_0 t),$$

the equilibrium corresponds to $A = 0$, $\Gamma_j = 1$, $\sum_j \exp(-in\tilde{\psi}_j) = 0$ ($n = 1, 2, \dots$).

Perturbing Eqs. (1)-(7), defining collective variables as in Ref. 12, and assuming a temporal dependence of the form $\exp(-i\Gamma t)$, the following dispersion relation is obtained

$$\Gamma^3 + \left[\dot{\psi}_0 - 2c^2(1-i\alpha)/\omega r_s^2 \right] \Gamma^2 + \frac{v(a_w f_B)^2}{N\gamma_0^3} \left(\frac{c}{r_s} \right)^2 \sigma \Gamma - 2ck_w \left(\frac{\gamma_r}{\gamma_0} \right)^2 \frac{v(a_w f_B)^2}{N\gamma_0^3} \left(\frac{c}{r_s} \right)^2 \sigma = 0,$$

where $\sigma = 2 \sum_j (1 - r_j^2/r_s^2) \exp(-2r_j^2/r_s^2)$. Perturbing B [defined by Eq. (8)], and numerically solving Eqs. (4) and (5) along with the cubic in Γ , one obtains the growth rate, efficiency, spot size, and α in the exponential regime for any given angular frequency ω .

Now, as is well-known the FEL interaction has the rather important property that the optical field tends to be guided by the electron beam. In an amplifier operating in the exponential regime, it is found that,

irrespective of the initial spot size, the radiation beam asymptotes to a unique spot size r_s and wavefront curvature ($\sim \alpha^{-1}$) (Refs. 8 and 9).

Figure 8 shows the growth rate Γ , efficiency η , matched spot size r_s and matched α in the case of a small input power as a function of $\delta = (\lambda/\lambda_{res}) - 1$, where $\lambda = 2\pi c/\omega$. The crosses are the results of single-bucket simulations. The curves are obtained from the linear stability analysis of the preceeding paragraph. It is seen that the agreement is quite good.

An important feature of Fig. 1 may now be understood with reference to Fig. 8. In the small signal regime - and therefore prior to particle trapping - Eqs. (1) - (7) may be linearized to show that there is no coupling between the various spectral components. In other words, the development of the spectral components proceeds independently and at a rate approximately equal to that indicated in Fig. 8. Referring to Fig. 1, it is thus seen that the larger growth rate of the sideband at $\delta = 0.77 \times 10^{-2}$ as compared to the carrier at $\delta = 2.51 \times 10^{-2}$ is consistent with Fig. 8.

It is also possible to set up and analyze a simple model of the FEL interaction so as to obtain an understanding of the gross features of the sideband instability ensuing from the synchrotron oscillation of the electrons in a large-amplitude carrier wave.

Neglecting diffraction and considering the fundamental mode of the optical field only, writing $\gamma_j = \gamma_r + \Delta\gamma_j(t)$, $\psi_j = \psi_{j0} + \Delta\psi_j(t)$, $a = |a^{(0)}| \exp[i\phi^{(0)}]$, the equilibrium is described by

$$\psi_{j0} + \phi^{(0)} + \alpha r_j^2 / r_s^2 = 2n_j\pi \quad , \quad (n_j \text{ is an integer})$$

$$\frac{d^2}{dt^2} \Delta\psi_j \approx -\Omega_j^2 \Delta\psi_j \quad , \quad (9)$$

$$|a^{(0)}| = \text{constant}.$$

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z}\right) \phi^{(0)} = \frac{v}{N\omega |a^{(0)}|} \left(\frac{2c}{r_s}\right)^2 \sum_j \frac{f_B a_w}{\gamma_r} \varepsilon_j, \quad (10)$$

where the synchrotron frequency is given by

$$\Omega_j = \left[\frac{ck_w \omega a_w f_B}{\gamma_r^2} |a^{(0)}| \varepsilon_j \right]^{1/2},$$

$\varepsilon_j = \exp(-r_j^2/r_s^2)$, and γ_r is defined by Eq. (3). Note that Eq. (10) describes a linear increase in the equilibrium phase. Inserting numerical values corresponding to the wiggler entrance, Eq. (10) yields a rate of increase that is about a factor of 4 greater than that observed in Fig. 3 (beyond 9 m). Due to diffraction, however, the spot size r_s increases along the wiggler. Thus, inserting an average value for r_s (between 9 m and 25 m) the rate obtained from Eq. (10) is within several tens of percent of the numerical value.

Next, perturbing the equilibrium state and assuming a dependence of the form $f^{(1)} \sim \exp[-i(\Delta\omega t - \Delta k z)]$ for the perturbation, one obtains the following dispersion relation (Appendix A)

$$1 = \frac{1}{X^2} - \frac{\left(2\beta ck_w / \Omega_B\right)^2}{(X-\zeta)^2} - \frac{\beta}{X(X-\zeta)}, \quad (11)$$

where

$$X = \Delta\omega \left(1 + a_w^2/2\right) / 2\Omega_B \gamma_r^2, \quad \zeta = c\Delta k \left(1 + a_w^2/2\right) / 2\Omega_B \gamma_r^2$$

$$\beta = \frac{a_w f_B}{2 |a^{(0)}| \gamma_r^5} \left(\frac{1 + a_w^2/2}{k_w r_s} \right)^2 \left(\sum_j \varepsilon_j / N \right),$$

and

$$\Omega_B^2 = \sum_j \Omega_j^2 / N. \quad (12)$$

The dispersion relation is obtained in the diagonal approximation, wherein terms of the form $f^{(1)} \Delta \psi_j$, $f^{(1)} \Delta \gamma_j$ are neglected. The dispersion relation in Eq. (11) is similar to that given in Ref. 6. The last term on the right-hand side of Eq. (11), which is due to terms proportional to $\gamma_j^{(1)}$ and was neglected in Ref. 6, increases the growth rate by about 10%. In reference to the electron distribution function shown in Fig. 4, it is interesting to note the similarity of Eq. (11) to the dispersion relation for a multi-stream system. Neglecting the third term on the right-hand side of Eq. (11), it is seen that - as discussed in Ref. 6 - the dispersion relation has a symmetrical form about the carrier ($\Delta k = 0$), with identical growth rate for $(-\Delta\omega, -\Delta k)$ and $(\Delta\omega, \Delta k)$. The spectrum in Fig. 5 is in approximate agreement with this general feature of Eq. (11).

Inserting numerical values into Eq. (11), one finds that the range of unstable modes encompasses that obtained in the simulations. The maximum growth rate obtained is about a factor of 5 larger than the average value observed for one of the fastest growing modes in the simulations [cf. Fig. 2]. That the growth rate deduced from Eq. (11) should exceed that observed in the simulations and the discrepancy in the rate of increase of the phase of the carrier are to be expected for two reasons. First, diffraction of the optical field [neglected in deriving Eq. (11)] is bound to reduce growth rates. Second, in the simulations the electrons are distributed throughout the ponderomotive bucket with synchrotron frequencies ranging from Ω_j down to zero, whereas the analysis leading to Eq. (11) assumes all electrons to bounce at the bottom of the ponderomotive wells, at the largest synchrotron frequency [cf. Eq. (9)].

Turning next to the case of the tapered magnet, the question arises as to why the growth rate of the sidebands is about a factor of 6 smaller than that in the untapered magnet. [cf. Figs. 2 and 6]. As is well-known, upon tapering, the electrons separate into roughly two groups. For the tapering employed, the decelerating group - which is responsible for stimulated emission - comprises 25-30% of the total number of electrons. However, the sideband growth rate for such a fraction of the number of electrons is still large compared to the observed value. There appears to be additional reasons for the small growth rate of the sidebands. In the context of a high-extraction FEL with a tapered magnet it is generally assumed that the change of parameters (a_w and/or k_w) is slow enough so that the action $I_j \equiv \int d\psi_j (\gamma_j - \gamma_r)$ is an adiabatic invariant.¹ For the case considered here, however, it turns out that $d\gamma_r/dt \approx \Omega_j(\gamma_j - \gamma_r)$, implying that the action I_j is not an invariant (Appendix B). Additionally, the γ_j of trapped electrons decrease more or less monotonically with time, modifying the "equilibrium" electron distribution on the same time scale as that of the synchrotron motion. Physically, as the radiation pulse slips relative to the electrons, it is modulated not only by the synchrotron motion but also by the temporal variation of γ_r . This has a detuning effect and renders the sidebands quasi-stable. Indeed, by increasing the rate of tapering and therefore the rate of variation of γ_r , the modulation of the optical pulse is reduced further, although the efficiency is also diminished due to increased particle detrapping.

Finally, it is important to remark that the tapering employed here leads to the same efficiency for a single-bucket, single-frequency case as for a many-bucket, multi-frequency case. In other words, for the given analytical form for $a_w(z)$, the optimal rate of tapering naturally leads to the observed reduction in the growth rate of sideband modes.

V. Conclusion

The work presented here contrasts the development of sideband frequencies in an untapered and in a tapered magnet. For the tapering employed, optimal operation (i.e., maximum efficiency) is achieved when the electrons trapped in the ponderomotive buckets are non-adiabatically decelerated, leading to substantial suppression of the sideband frequencies relative to the untapered system.

Acknowledgement

Discussions with Dr. W. P. Marable are gratefully acknowledged. This work was supported by SDIO and managed by SDC.

Appendix A: Multi-frequency linear stability analysis

In this appendix some of the details leading to the dispersion relation in Eq. (11) are presented.

Upon writing $\gamma_j = \gamma_r + \Delta\gamma_j + \gamma_j^{(1)}$, $\psi_j = \psi_{j0} + \Delta\psi_j + \psi_j^{(1)}$, $\phi = \phi^{(0)} + \phi^{(1)}$, $|a| = |a^{(0)}| + |a^{(1)}|$, and neglecting terms comprising products of the form $\Delta\gamma_j \psi_j^{(1)}$, $\Delta\psi_j \gamma_j^{(1)}$, etc., one obtains, at first order,

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z}\right) |a^{(1)}| \approx \frac{v}{N\omega} \left(\frac{2c}{r_s}\right)^2 \sum_j \frac{f_B a_w}{\gamma_r} \varepsilon_j \left(\psi_j^{(1)} + \phi^{(1)}\right),$$

$$|a^{(0)}| \left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z}\right) \phi^{(1)} \approx - \frac{v}{N\omega} \left(\frac{2c}{r_s}\right)^2 \sum_j \frac{f_B a_w}{\gamma_r} \varepsilon_j \left(\frac{|a^{(1)}|}{|a^{(0)}|} + \frac{\gamma_j^{(1)}}{\gamma_r}\right), \quad (A1)$$

$$\frac{d}{dt} \gamma_j^{(1)} \approx - \frac{\omega a_w f_B}{2\gamma_r} \varepsilon_j |a^{(0)}| \left(\psi_j^{(1)} + \phi^{(1)}\right),$$

$$\frac{d}{dt} \psi_j^{(1)} = 2 c k_w \gamma_j^{(1)} / \gamma_r,$$

Note that the term proportional to $\gamma_j^{(1)}$ on the right-hand side of Eq. (A1) was neglected in Ref. 6 where a similar analysis of sideband growth was presented.

Defining collective variables

$$x_k = \sum_j \varepsilon_j^k \psi_j^{(1)} / N, \quad y_k = \sum_j \varepsilon_j^k \gamma_j^{(1)} / N, \quad E_k = \sum_j \varepsilon_j^k / N,$$

one obtains

$$\left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z}\right) |a^{(1)}| = \frac{v}{\omega} \left(\frac{2c}{r_s}\right)^2 \frac{f_B a_w}{\gamma_r} \left(x_1 + E_1 \phi^{(1)}\right),$$

$$|a^{(0)}| \left(\frac{\partial}{\partial t} + c \frac{\partial}{\partial z}\right) \phi^{(1)} = - \frac{v}{\omega} \left(\frac{2c}{r_s}\right)^2 \frac{f_B a_w}{\gamma_r} \left(\frac{y_1}{\gamma_r} + \frac{E_1 |a^{(1)}|}{|a^{(0)}|}\right),$$

$$\frac{d}{dt} X_1 = \frac{2ck_w}{\gamma_r} y_1 ,$$

$$\frac{d}{dt} y_1 = - \frac{\omega a_w f_B}{2\gamma_r} |a^{(0)}| \left(X_2 + E_2 \phi^{(1)} \right) ,$$

$$\frac{d}{dt} X_2 = \frac{2ck_w}{\gamma_r} y_2 ,$$

etc.

This hierarchy may be truncated by assuming $X_2 = \Sigma \epsilon_j X_1 / N$, $E_2 = \Sigma \epsilon_j E_1 / N$,
whence

$$\frac{d^2}{dt^2} X_1 = - \Omega_B^2 \left(X_1 + E_1 \phi^{(1)} \right) ,$$

$$\frac{d}{dt} y_1 = - \frac{\gamma_r \Omega_B^2}{2ck_w} \left(X_1 + E_1 \phi^{(1)} \right) ,$$

where Ω_B is defined by Eq. (12).

Assuming perturbations of the form $\exp [-i(\Delta\omega t - \Delta kz)]$, one obtains the dispersion relation in Eq. (11), where the third term on the right-hand side is due to the term proportional to $\gamma_j^{(1)}$ on the right-hand side of Eq. (A1).

Appendix B: Temporal variation of adiabatic invariant

In this appendix, the law of variation of the action variable for the tapered magnet is obtained.

For a synchronous particle, the energy and phase evolve according to

$$\frac{d}{dt} \gamma_r = - \frac{\omega a_w f_B}{2\gamma_r} |a| \varepsilon_r \sin(\psi_r + \phi + \alpha r_r^2 / r_s^2) ,$$

$$\frac{d}{dt} \psi_r = 0 .$$

Defining $\gamma_j = \gamma_r + \delta\gamma_j$, $\psi_j = \psi_r + \delta\psi_j$, the Hamiltonian for small amplitude oscillations is

$$H = \frac{ck_w}{\gamma_r} \delta\gamma_j^2 + \frac{\omega a_w f_B |a|}{4\gamma_r} \varepsilon_r \cos(\psi_r + \phi + \alpha r_r^2 / r_s^2) \delta\psi_j^2 .$$

Defining the synchrotron frequency of such a particle,

$$\Omega_{\text{syn}} = \left[ck_w \omega a_w f_B |a| \varepsilon_r \cos(\psi_r + \phi + \alpha r_r^2 / r_s^2) \right]^{1/2} / \gamma_r ,$$

following Ref. 13, the action (I_j) angle (w_j) variables are found to evolve according to

$$\frac{d}{dt} I_j = - \frac{I_j}{\gamma_r} \frac{d\gamma_r}{dt} \cos 2w_j ,$$

$$\frac{d}{dt} w_j = \Omega_{\text{syn}} + \frac{1}{2\gamma_r} \frac{d\gamma_r}{dt} \sin 2w_j .$$

These formulae are valid for the length of the magnet lying between 6 m and 21 m where Ω_{syn} is found to be remarkably constant in the simulations.

Thus the variation of γ_r is the only factor contributing to the breaking of the adiabatic invariant I_j . (Generalization of the formulae to include the variation of Ω_{syn} in the other sections of the magnet is straightforward.)

The change in the action variable over a synchrotron period is then found to be given by

$$\frac{|\Delta I|}{I} \lesssim \pi \left(\frac{d\gamma_r/dt}{\gamma_r \Omega_{\text{syn}}} \right)^2 ,$$

$$\sim 10^{-3} .$$

It is important to notice that although this variation is relatively small, it is nevertheless much faster than the usual case for an adiabatic invariant where $\Delta I \rightarrow 0$ exponentially as $d\gamma_r/dt \rightarrow 0$.

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Table I. Parameters of the Paladin Experiment

Electron Beam

Current	2 kA
Energy	50 MeV
Radius	0.45 cm

Magnet

Induction	2.3 kG
Period	8 cm
Length	25 m

Radiation Field

Wavelength	10.6 μm
Initial Spot Size	0.36 cm

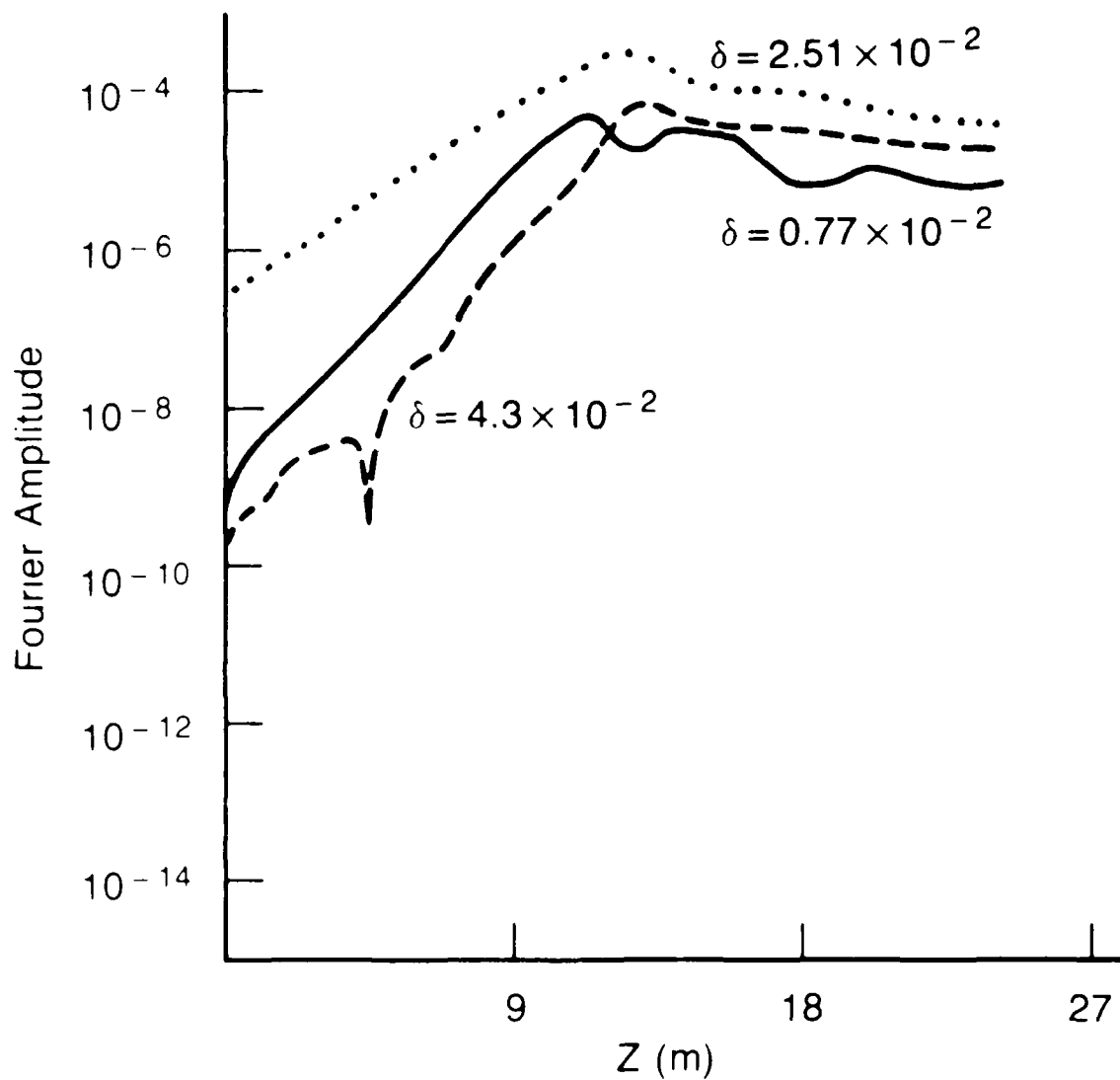


Fig. 1 Evolution of 3 spectral components along the wiggler when 100 W is input into the 10.6 μm wavelength ($\delta = 2.51 \times 10^{-2}$) at the entrance to the wiggler.

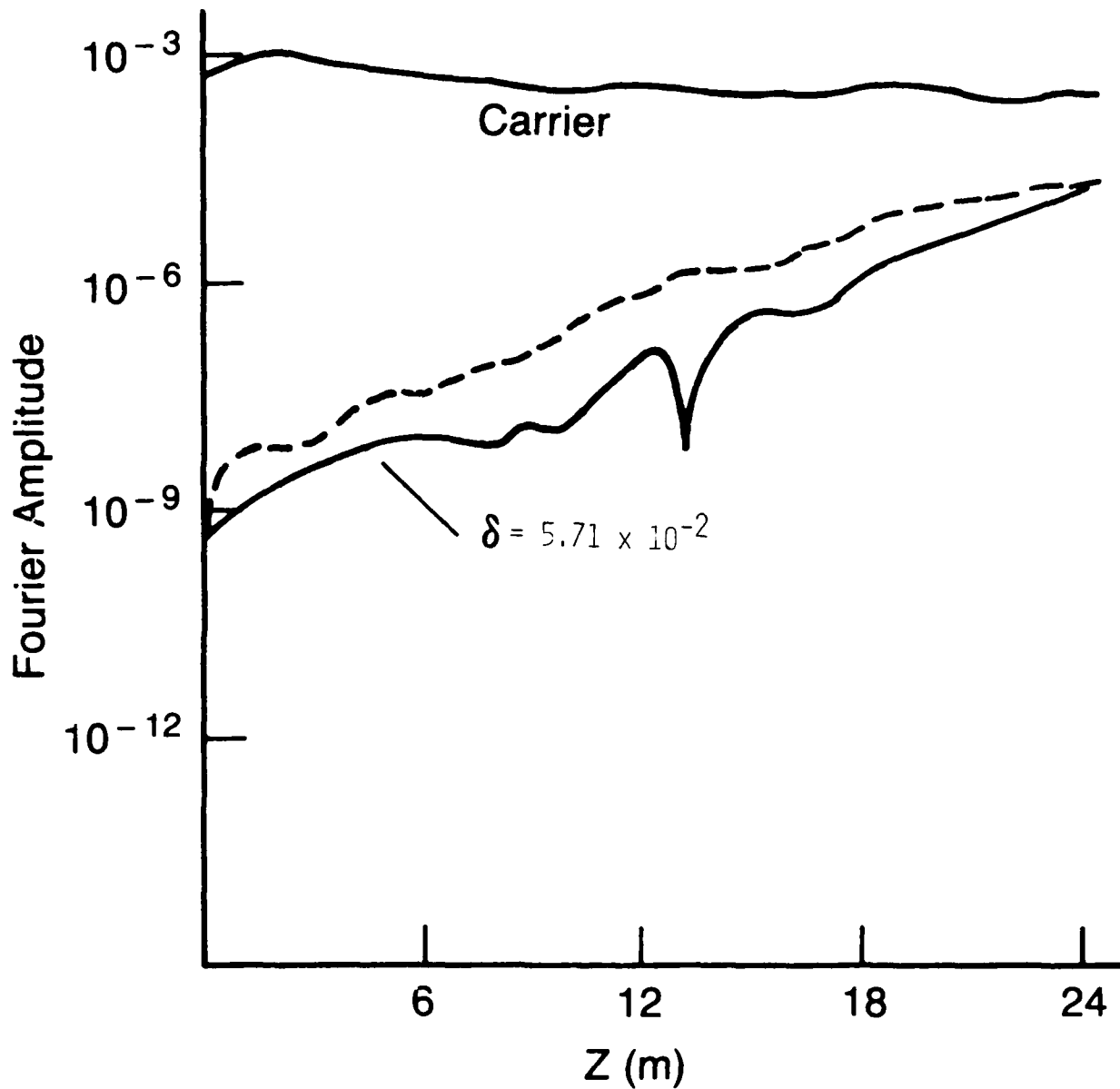


Fig. 2 Evolution of carrier ($10.6 \mu\text{m}$) starting from 800 MW in an untapered wiggler. The dashed curve indicates the upper bound for (or the envelope of) the rest of the spectrum. Also shown is the evolution of the sideband at $\delta = 5.71 \times 10^{-2}$ ($10.93 \mu\text{m}$).

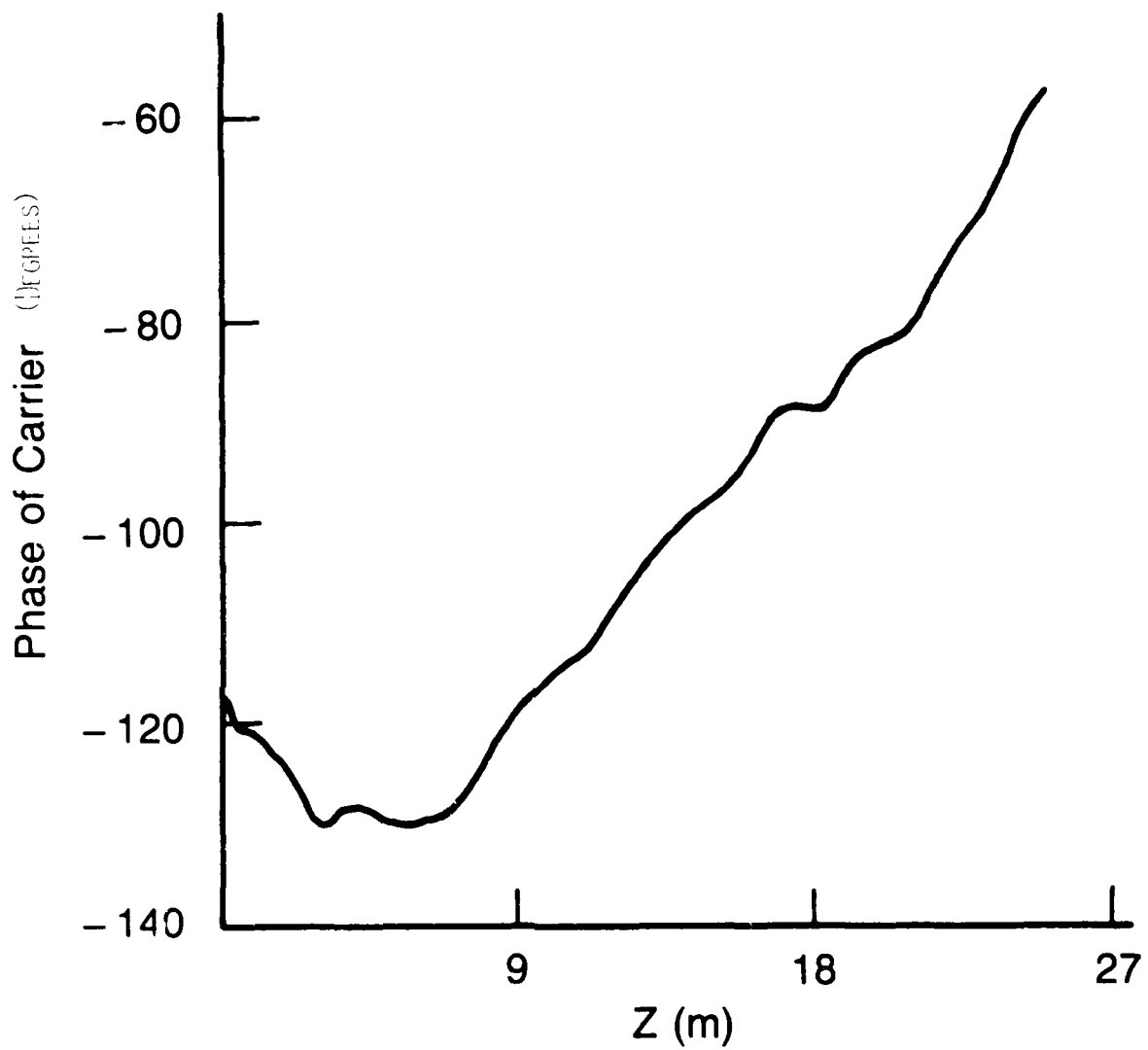


Fig. 3 Development of the phase of the carrier (in degrees) along the wiggler.

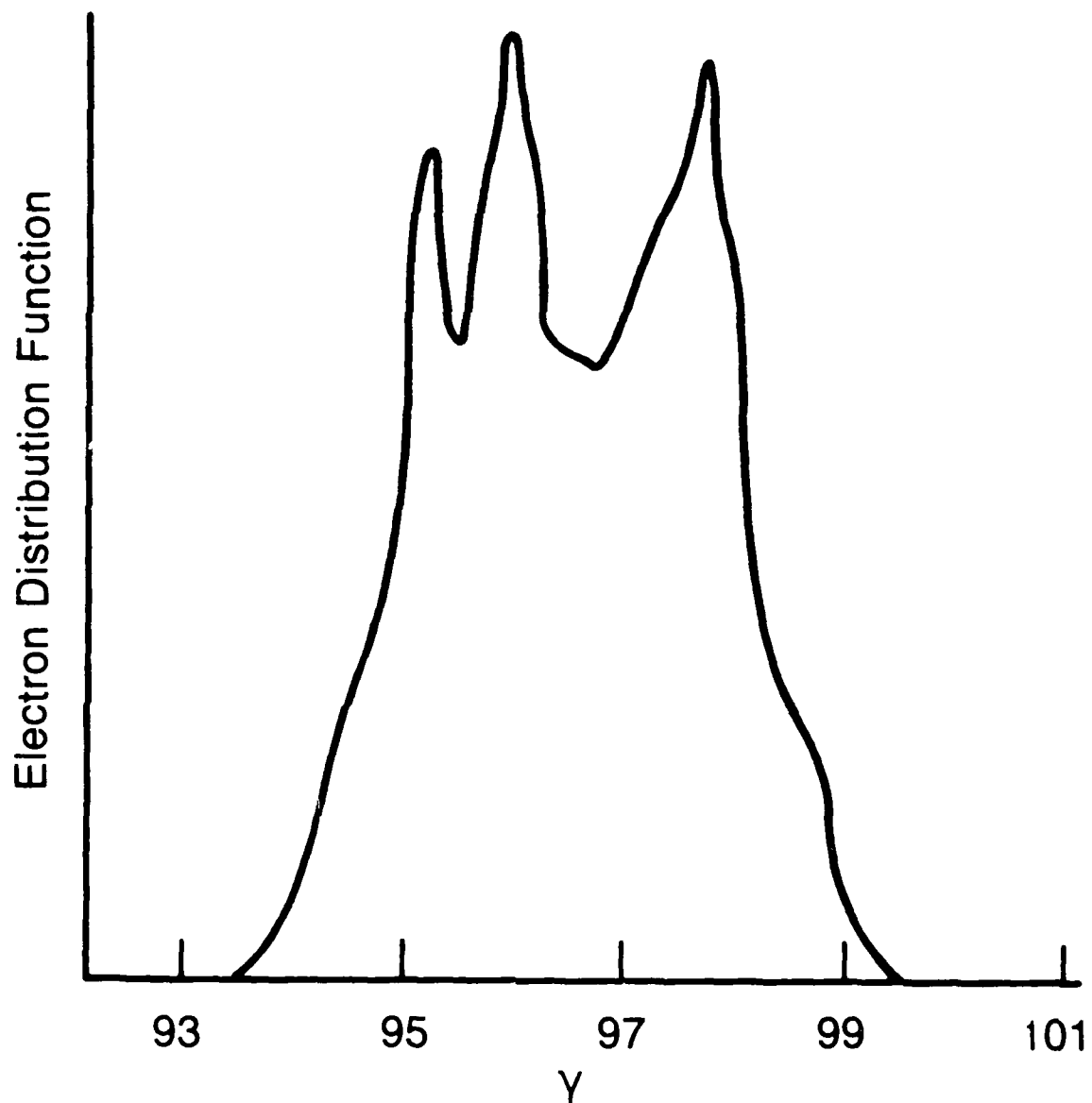


Fig. 4 Electron distribution function (i.e., number of electrons) versus relativistic mass factor γ at the end of the wiggler. (Ordinate scale is linear).

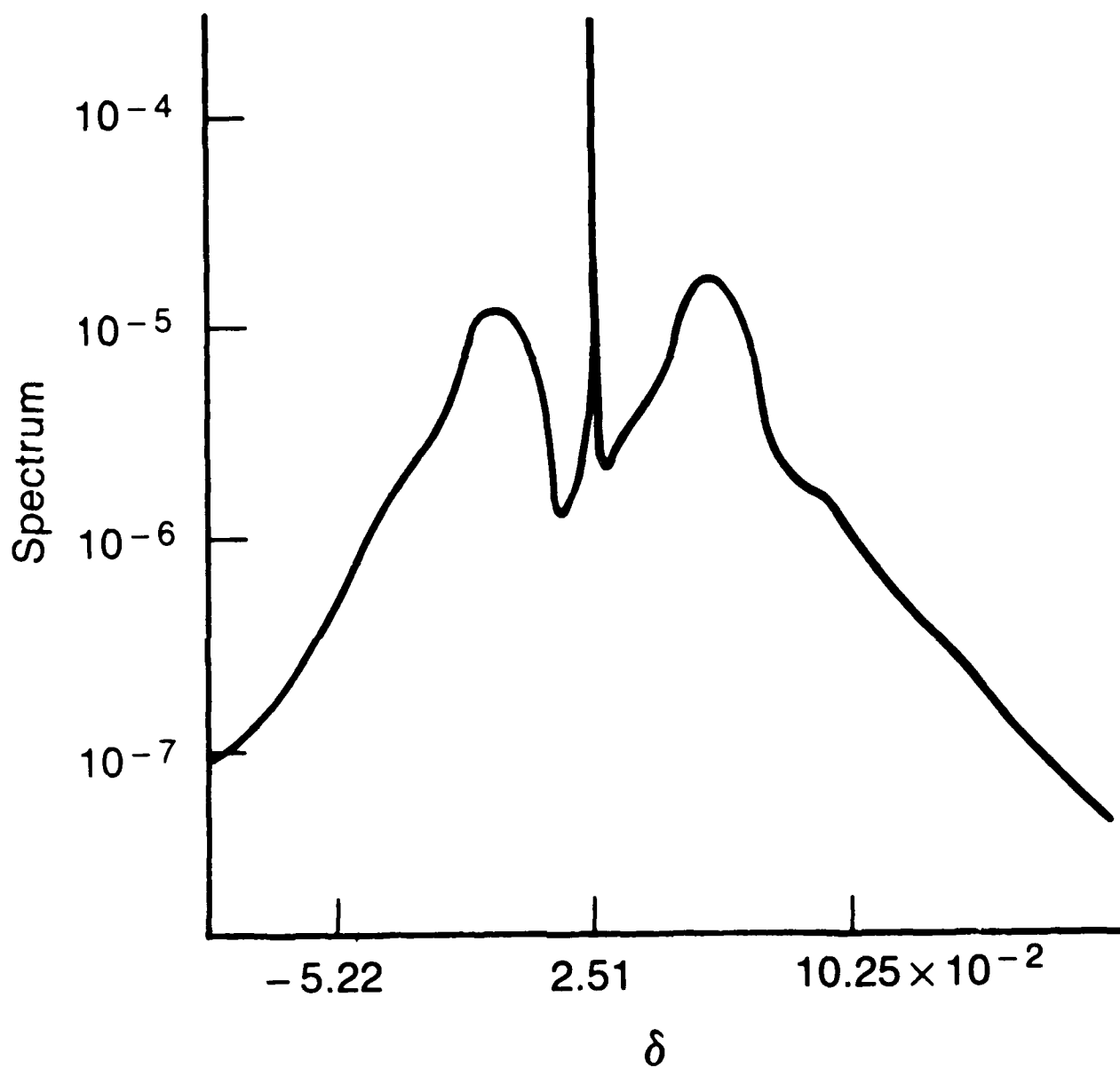


Fig. 5 Spectrum of optical field versus $\delta \equiv \lambda/\lambda_{\text{res}} - 1$ at the end of the wiggler.

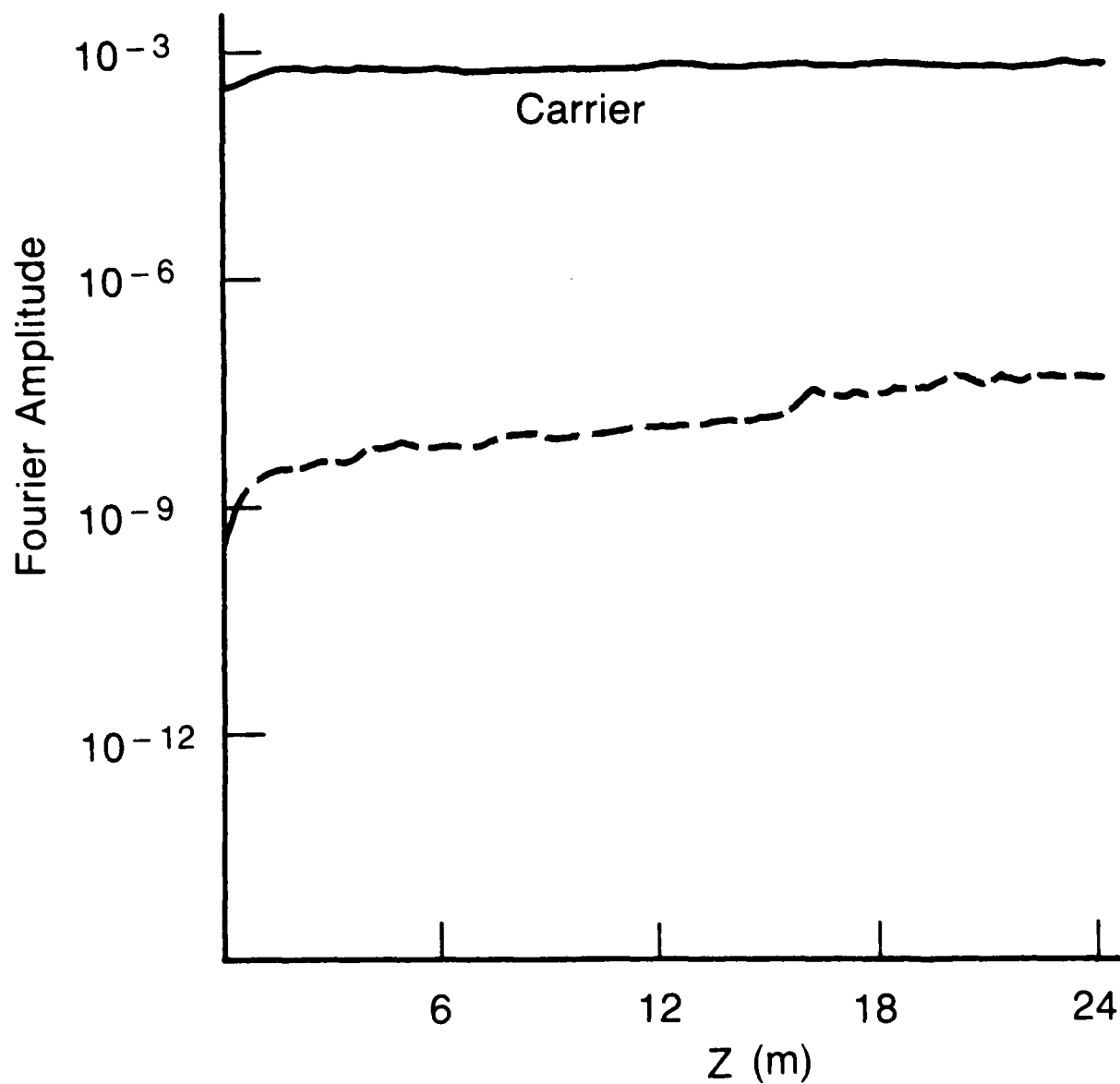


Fig. 6 Evolution of carrier ($10.6 \mu\text{m}$) starting from 800 MW in a tapered wiggler. The dashed curve indicates the upper bound for the rest of the spectrum.

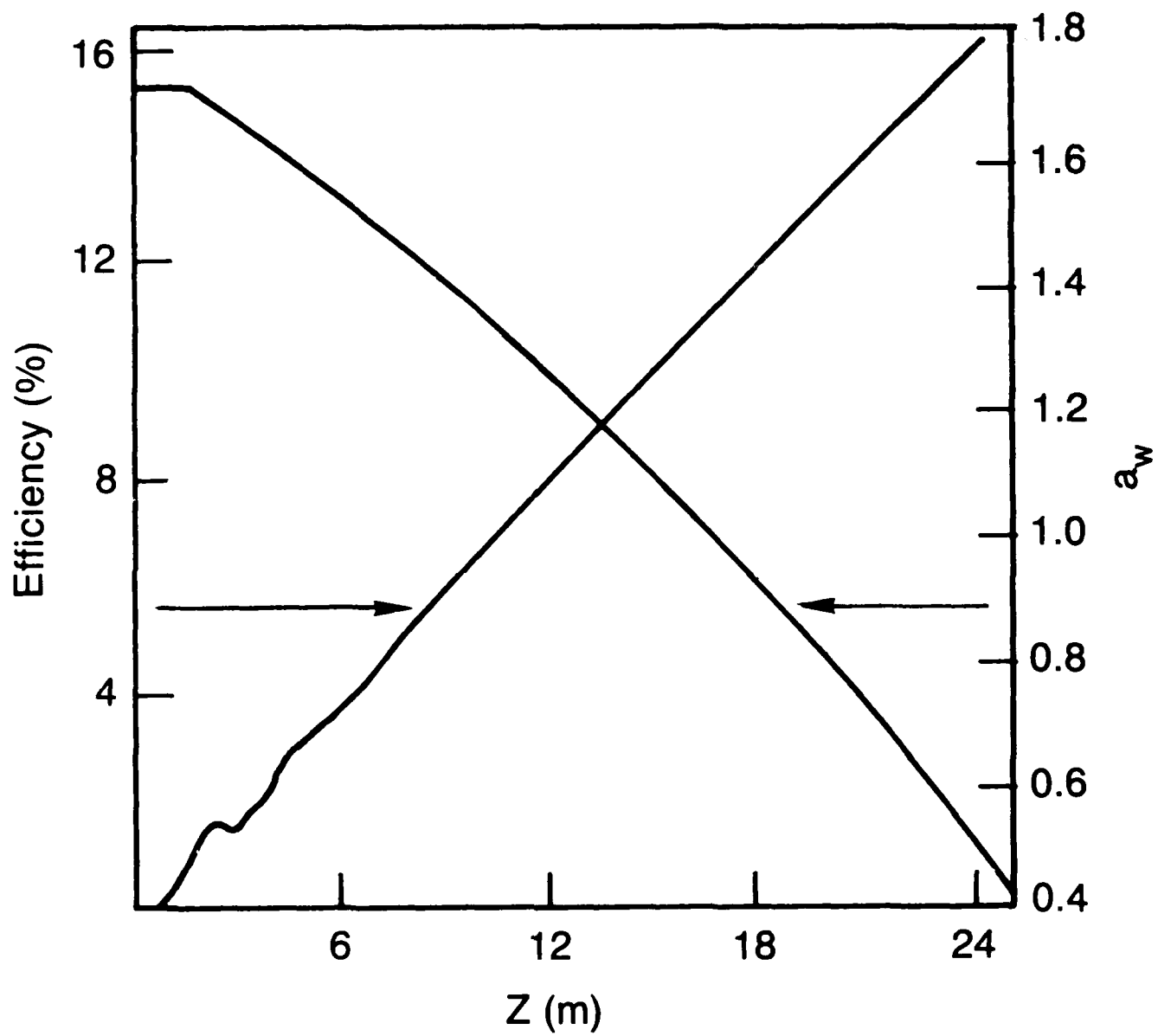


Fig. 7 Efficiency (%) and a_w along the length of the tapered magnet.

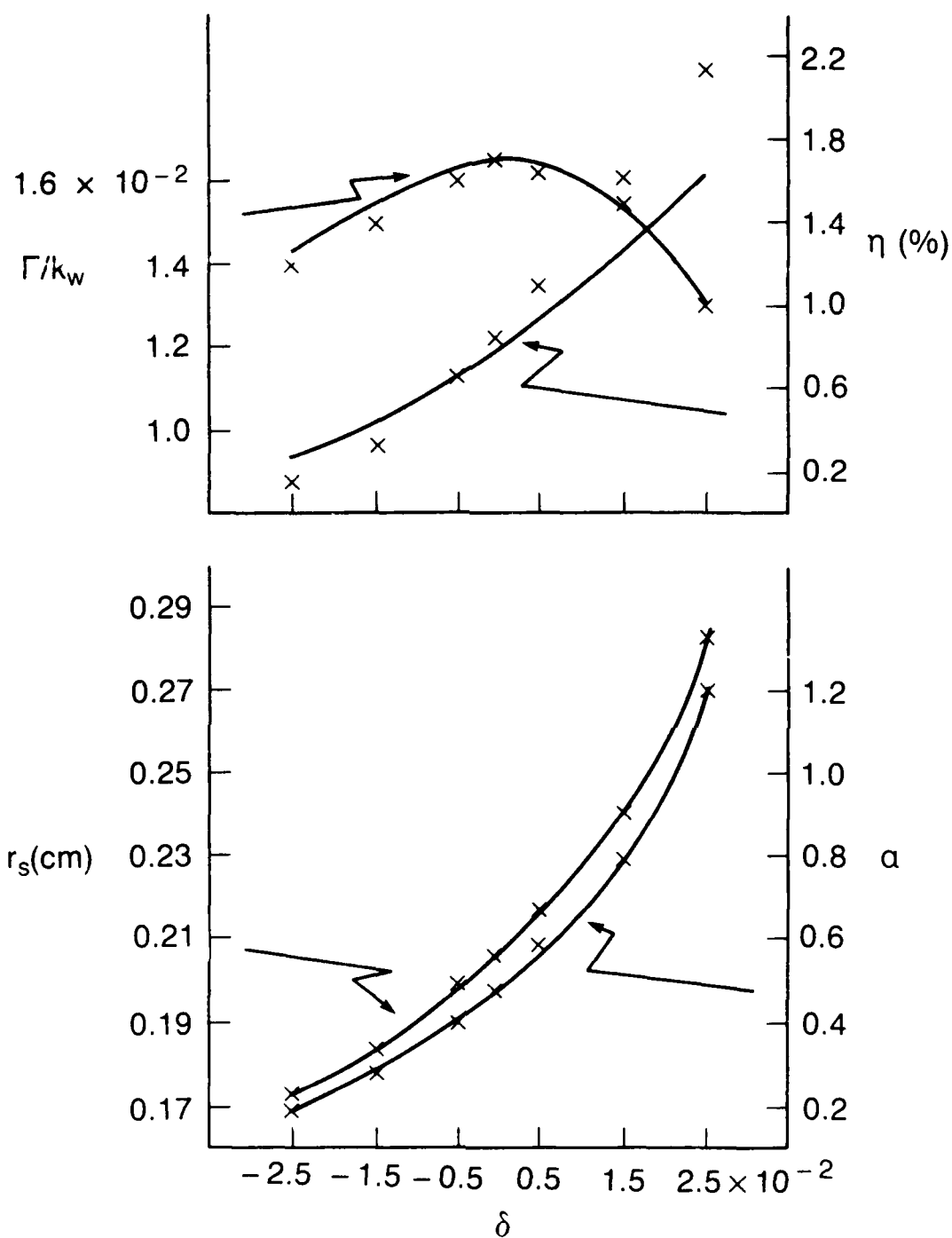


Fig. 8 Normalized growth rate Γ/k_w , efficiency $\eta(\%)$, spot size r_s , and α versus $\delta \equiv (\lambda/\lambda_{\text{res}})-1$, where $\lambda_{\text{res}} = 10.34 \mu\text{m}$. Crosses represent results of simulations. Curves are obtained from the linearized equations for the exponential, matched-optical-field regime. The electron beam radius is 0.3 cm .

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